

Photon polarization as a probe for quark-gluon plasma dynamics

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Abstract

Prospects of measuring polarized photons emitted from a quark-gluon plasma (QGP) are discussed. In particular, the detection of a possible quark spin polarization in a QGP using circularly polarized photons emitted from the plasma is studied. Photons leave the QGP without further interaction and thus provide a primary probe for quark polarization within the QGP. We find that photon polarization cannot solely arise due to a possible QGP momentum space anisotropy, but may be enhanced due to it. In particular, for oblate momentum distributions and high photon energies, quark polarization is efficiently transferred to photon polarization. The role of competing sources of polarized photons in heavy-ion collisions is discussed.

Key words: Quark gluon plasma

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The discovery of the quark-gluon plasma (QGP), a new state of matter at temperatures above $T \gtrsim 170$ MeV, in heavy-ion collisions at CERN/SPS [1] and RHIC [2] leaves open many questions about its properties and its internal structure: If the naive picture of free quarks and gluons is obsolete, is the QGP rather a strongly coupled system [3], or more adequately described as a fluid with low viscosity [4]? How does the QGP apparently thermalize so fast ($\lesssim 1$ fm/c) [5]? Are quarks in the QGP polarized [6]? In order to learn more about the QGP, one would like to have a primary probe directly from its interior, which is difficult due to its very short lifetime $\tau \sim 5$ fm/c. Such a probe is given by direct thermal photons that, when produced, typically escape the plasma without further interaction. The PHENIX experiment at RHIC has confirmed the good agreement between the photon rate measured and perturbative calculations [7]. It has been suggested that direct photons

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could provide information about a possible momentum anisotropy within the QGP [8]. Polarized photon emission from the QGP has been previously studied as a geometrical surface effect [9], but to our knowledge not from the bulk.

Global quark polarization within the QGP has been proposed by Liang and Wang [6] to occur in non-central heavy-ion collisions by a simple mechanism: Due to a finite impact parameter, the bulk of the QGP is kicked to rotate within the reaction plane, and the resulting potentially large orbital angular momentum may lead to quark spin polarization due to spin-orbit coupling (see also [10,11]). A possible transfer of quark polarization on massive particles was studied for the case of Λ -hyperons. Experimentally, an upper limit of the polarization $|P_{\Lambda,\bar{\Lambda}}| \lesssim 0.02$ has been obtained at RHIC [12,13], while recent calculations suggest $P_{\Lambda} \approx -0.05$ [14,15]. It should be noted though that hyperon polarization is affected by all stages of the collision, including the later hadronization and the interacting hadron gas phase, to an unknown degree [16].

In this Letter, we show that global QGP polarization would effectively lead to a polarization of photons. Photons are a primary probe as they are likely to leave the plasma without further interaction. Both angular distribution and polarization of the emitted photons depend on the quark polarization. Particularly, we point out the possibility to detect a possible global quark polarization using circularly polarized photons. We also show that a momentum-anisotropy alone does not lead to a polarization of the emitted photons. Therefore, the polarization of photons can be facilitated as a direct signal of the quark spin polarization.

Thermal photons in the QGP can be produced through Compton scattering of (anti-)quarks and gluons, $qg \rightarrow q\gamma$ ($\bar{q}g \rightarrow \bar{q}\gamma$), and annihilation of quarks and antiquarks, $q\bar{q} \rightarrow g\gamma$. The bulk of the QGP is dominated by up and down quarks, for which the ultrarelativistic limit $T \gg m$ for the current quark masses m applies. In thermal equilibrium, photons are emitted isotropically, thermally distributed, and unpolarized. In a heavy-ion collision, however, the QGP expands along the beam axis which induces momentum space anisotropy, and the produced photons may reflect this in their rapidity dependence [8]. For clarity, we will use the term “anisotropy” for momentum space anisotropy within the QGP, “polarization” for quark spin or photon polarization, and we use the temperature T as hard energy scale ($q_{\text{hard}} = T$), even though it is properly defined only in the isotropic case. Also, $\hbar = c = k_B = 1$.

The spin of a quark with momentum $\kappa = (\kappa_0, \boldsymbol{\kappa})$ can be defined in its rest frame [17]. Under a boost along $\boldsymbol{\kappa}$, only the longitudinal component of the spin four-vector $s_{\text{rest}} = (0, \mathbf{s}_{\text{rest}})$ changes, the transverse contribution being invariant. For an ensemble of quarks with momentum $\boldsymbol{\kappa}$, the average spin \mathbf{s}_{rest} can be split according to $\mathbf{s}_{\text{rest}} = p^{\parallel} \hat{\boldsymbol{\kappa}} + p^{\perp} \hat{\mathbf{s}}^{\perp}$ [18], where $p^{\parallel} = p^{\parallel}(\hat{\boldsymbol{\kappa}})$ and

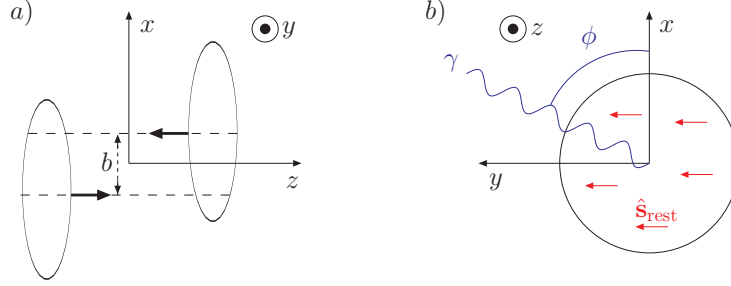


Figure 1. (Color online). (a) Heavy ion collisions with non-zero impact parameter b may produce a global spin polarization of quarks along $\hat{\mathbf{y}}$ as proposed in Ref. [6]. (b) Such a global quark spin $\hat{\mathbf{s}}_{\text{rest}} = \hat{\mathbf{y}}$ could be detected through the azimuthal ϕ -dependence of photons emitted from the plasma.

$p^\perp = p^\perp(\hat{\boldsymbol{\kappa}})$ denote the degree of polarization along the propagation direction $\hat{\boldsymbol{\kappa}} \equiv \boldsymbol{\kappa}/|\boldsymbol{\kappa}|$ or along a transverse direction $\hat{\mathbf{s}}^\perp = \hat{\mathbf{s}}^\perp(\hat{\boldsymbol{\kappa}})$ with $\hat{\boldsymbol{\kappa}} \cdot \hat{\mathbf{s}}^\perp = 0$. States with partial polarization can be expressed in the ultrarelativistic limit using the density matrices $\rho_\pm = \not{\epsilon} P_\pm$ with $P_\pm = (1 \mp p^\parallel \gamma^5 + p^\perp \gamma^5 \not{\epsilon}^\perp)/2$, where the upper (lower) sign refers to particles (antiparticles) [17]. Partial polarization $|\mathbf{s}_{\text{rest}}|^2 < 1$ is understood as a linear combination $p^\parallel = p^\parallel(\uparrow) - p^\parallel(\downarrow)$ with $p^\parallel(\uparrow) + p^\parallel(\downarrow) = 1$ (and likewise for p^\perp) of fully polarized contributions $|\mathbf{s}_{\text{rest}}|^2 = 1$, for which $P_\pm^2 = P_\pm$ are projectors [17].

In principle, for a certain point in coordinate space, momentum dependent polarization functions $p^\parallel(\kappa)$, $p^\perp(\kappa)$, and $\hat{\mathbf{s}}^\perp(\kappa)$ for each (anti-)quark species would completely specify the quark polarization content within the QGP. Ideally, these functions should follow from some first-principle calculation in the medium. Since they are unknown, we proceed in a pragmatic way: in accordance with global polarization proposed for non-central collisions [6,15] (see Fig. 1), we assume that the spin of each quark in its rest frame is aligned along the same direction $\hat{\mathbf{s}}_{\text{rest}}(\kappa) = \hat{\mathbf{y}}$ and that particles with same energy κ_0 share the same degree of polarization $p_{\text{rest}}(\kappa_0) = |\mathbf{s}_{\text{rest}}(\kappa_0, \boldsymbol{\kappa})|$. Since the main polarization effect is expected to arise from hard modes with energy $\kappa_0 \sim T$ [14,15], we use the following model for the energy dependence of $p_{\text{rest}}(\kappa_0)$:

$$p_{\text{rest}}(\kappa_0) = \bar{p}_{\text{rest}} \Theta(\kappa_0 - k^*). \quad (1)$$

The energy threshold k^* is chosen equal to the separation scale between soft and hard modes, to be specified below Eq. (6). Modes that greatly exceed the temperature are automatically cut off by the thermal distribution functions of the plasma.

The calculation of the production rate of photons with momentum $q = (E, \mathbf{q})$ requires a separation into hard ($\sim T$) and soft ($\sim gT$) momentum transfer contributions [19,8]. We restrict ourselves to hard $2 \leftrightarrow 2$ particle processes and corresponding soft processes, and will not consider processes involving

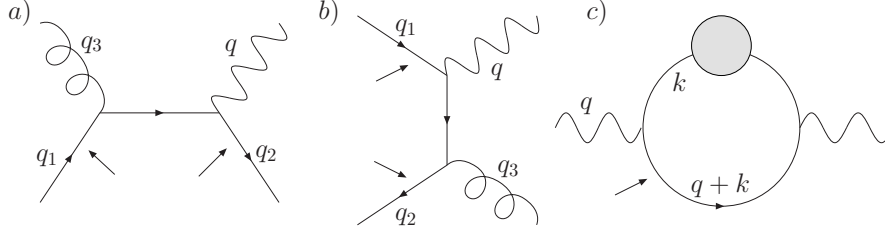


Figure 2. Feynman diagrams corresponding to (a) Compton scattering, (b) quark-antiquark annihilation, and (c) photon self-energy (straight, curly, and wiggly lines represent quarks, gluons, and photons). Arrows indicate where polarized states are used in the calculation. The gray filled circle indicates a soft propagator dressed by the medium.

more particles, like bremsstrahlung and inelastic pair annihilation. A consistent treatment of the latter processes requires an analysis of multiple soft scattering [20], which in the anisotropic case still poses fundamental open questions due to the possibility of plasma instabilities [8,5]. In the isotropic case, these processes change the total photon rate in the relevant energy range $2.5 \leq E/T \leq 10$ by a factor of two [20]. For hard momenta, standard perturbation theory is used to calculate the rate contributions from Compton scattering R_c or annihilation R_a

$$E \frac{dR_c}{d^3q} = 20\pi \int_{\delta} f_{\xi}^F(\mathbf{q}_1) f_{\xi}^B(\mathbf{q}_3) \left(|\mathcal{M}_c^S|^2 - f_{\xi}^F(\mathbf{q}_2) |\mathcal{M}_c|^2 \right),$$

$$E \frac{dR_a}{d^3q} = \frac{320\pi}{3} \int_{\delta} f_{\xi}^F(\mathbf{q}_1) f_{\xi}^F(\mathbf{q}_2) \left(1 + f_{\xi}^B(\mathbf{q}_3) \right) |\mathcal{M}_a|^2, \quad (2)$$

with $\int_{\delta} \equiv \int d^3q_1 d^3q_2 d^3q_3 / (2\pi)^9 \delta^{(4)}(q_1 + q_2 - q_3 - q) / (8E_1 E_2 E_3)$ (the momenta are defined in Fig. 2), and fermionic or bosonic distribution functions $f_{\xi}^{F/B}$. The numerical prefactors include summing over $N_c = 3$ colors and the fractional charge factors of the considered up and down quarks that dominate the rate due to their low mass. Pauli blocking has to be taken into account properly by including the scattering Compton matrix element either polarized (\mathcal{M}_c), or summed over outgoing quark polarization (\mathcal{M}_c^S). For anisotropic momentum distributions, the isotropic distribution functions $f_{\text{iso}}^{F/B}$ are assumed to be stretched along the beam axis \mathbf{z} in momentum space [21], $f_{\xi}^{F/B}(\mathbf{k}) = N(\xi) f_{\text{iso}}^{F/B}(k_{\xi})$, with $k_{\xi}^2 = k_x^2 + k_y^2 + (1 + \xi)k_z^2$, the anisotropy parameter $\xi > -1$, and a normalization factor which we choose as $N(\xi) = 1$ as in [8]. A different normalization is discussed later.

Before looking at globally polarized quarks, we first want to study the polarization of emitted photons from a QGP without quark polarization. In the following, we present the result for the Compton scattering diagram Fig. 2(a), since the corresponding scattering matrix \mathcal{M}_a for the annihilation diagram

Fig. 2(b) follows from crossing symmetry. For polarized Compton scattering with electron charge e and strong coupling g we find

$$|\mathcal{M}_c|^2 = \frac{e^2 g^2}{6} \left(\frac{q \cdot q_1}{q \cdot q_2} + \frac{q \cdot q_2}{q \cdot q_1} - 2m^2 \left| \frac{\epsilon \cdot q_1}{q \cdot q_1} - \frac{\epsilon \cdot q_2}{q \cdot q_2} \right|^2 \right). \quad (3)$$

In this equation we have kept a finite quark mass m dependence to demonstrate how the polarization vector $\epsilon(q)$ would enter the cross section. This formula leads to the well-known Klein-Nishina differential cross section [22]. For isotropic momentum distributions any polarization effect averages out. However, for anisotropic momentum distributions, finite quark masses may give rise to linear polarization of photons, even if the quarks in the plasma are unpolarized. But in the ultrarelativistic limit considered, Eq. (3) trivially reduces to the unpolarized case. For the soft part, the dressed quarks in the medium acquire a thermal mass of the parametric order gT . If this thermal mass adopts the role of an effective quark mass, photon polarization in an anisotropic medium would not yet be excluded. To address this question, we calculate the soft contributions for unpolarized quarks and polarized photons. These can be obtained from the imaginary part of the photon self-energy, which corresponds to the 12-component of the photon self-energy Π in Fig. 2(c) in the real-time formalism [19]

$$E \frac{dR_{\text{soft}}}{d^3q} = -\frac{i}{2(2\pi)^3} \epsilon_\mu^* \epsilon_\nu \Pi_{12}^{\mu\nu}(q). \quad (4)$$

The unpolarized photon form [23,8] is recovered by applying the Ward identity. For general polarization vectors $\epsilon(q)$ one finds within the Hard Loop approximation (which assumes $gT \sim k \ll q \sim T$) [19]

$$i\epsilon_\mu^* \epsilon_\nu \Pi_{12}^{\mu\nu}(q) = e^2 \frac{20f_\xi^F(\mathbf{q})}{3|\mathbf{q}|} \int \frac{d^3k}{(2\pi)^3} \text{Im } q \cdot S_R^*(k) \quad (5)$$

with the retarded dressed fermionic propagator $S_R^*(k)$ [24]. This result is independent of the polarization vector and is just 1/2 of the unpolarized result [23,8]. This means that within our model, photons emitted from a QGP without quark polarization are unpolarized, even if the distribution functions show momentum anisotropy. In other words, the quarks themselves have to be polarized in order to observe photon polarization, and we can directly relate photon polarization to quark polarization.

We analyze the quark spin dependence by using polarized states ρ_\pm as indicated in Fig. 2, which extends Eq. (3) in the ultrarelativistic case to

$$\begin{aligned}
|\mathcal{M}_c|^2 = & \frac{e^2 g^2}{12} \left\{ \left(1 + p^\parallel(q_1)p^\parallel(q_2)\right) \left(\frac{q \cdot q_1}{q \cdot q_2} + \frac{q \cdot q_2}{q \cdot q_1}\right) \right. \\
& + \left(p^\parallel(q_1) + p^\parallel(q_2)\right) \left(\frac{q \cdot q_2}{q \cdot q_1} - \frac{q \cdot q_1}{q \cdot q_2}\right) i \det |\epsilon \epsilon^* \hat{\mathbf{q}}| \\
& - 2p^\perp(q_1)p^\perp(q_2) \left[\left(q \cdot \hat{s}_1^\perp\right) \left(q \cdot \hat{s}_2^\perp\right) \left(\frac{1}{q \cdot q_1} - \frac{1}{q \cdot q_2}\right) \right. \\
& \left. \left. + \hat{s}_1^\perp \cdot \hat{s}_2^\perp - \frac{(q \cdot \hat{s}_1^\perp)(q_1 \cdot \hat{s}_2^\perp)}{q \cdot q_1} - \frac{(q \cdot \hat{s}_2^\perp)(q_2 \cdot \hat{s}_1^\perp)}{q \cdot q_2} \right] \right\}
\end{aligned} \tag{6}$$

with $\hat{s}_1^\perp = \hat{s}^\perp(q_1)$ and $\hat{s}_2^\perp = \hat{s}^\perp(q_2)$. For unpolarized outgoing quarks one has $|\mathcal{M}_c^S|^2 = 2|\mathcal{M}_c|^2$, with $p^\parallel(q_2) = p^\perp(q_2) = 0$. The pieces quadratic in the polarization, proportional to $p^\parallel(q_1)p^\parallel(q_2)$ and $p^\perp(q_1)p^\perp(q_2)$, vanish independently in the infrared (IR) limit through the angular integrations in (2), but they may still contribute at larger momentum transfers as we find numerically. Only the piece linear in the polarization, proportional to $p^\parallel(q_1) + p^\parallel(q_2)$, survives in the IR limit. This piece has a matching ultraviolet contribution from the soft sector: Due to the assumption in Eq. (1) we only need to take into account polarized states for the hard quark line in Fig 2(c).

For polarized quarks and photons, Eq. (5) is extended to

$$\begin{aligned}
i\epsilon_\mu^* \epsilon_\nu \Pi_{12}^{\mu\nu}(q) = & \left(1 + p^\parallel(q) i \det |\epsilon \epsilon^* \hat{\mathbf{q}}|\right) \\
& \times e^2 \frac{10 f_\xi^F(\mathbf{q})}{3|\mathbf{q}|} \int \frac{d^3 k}{(2\pi)^3} \text{Im } q \cdot S_R^*(k).
\end{aligned} \tag{7}$$

The determinant in this expression also occurs in the hard part in Eq. (6). It vanishes for linear polarization, but contributes for right-(left)-handed circular polarization with $i \det |\epsilon \epsilon^* \hat{\mathbf{q}}| = +(-)1$. The photon rate EdR/d^3q is finally given by adding the results of Eq. (2) and (4), with an intermediate cutoff $k^* \sim \sqrt{g}T$ that is varied to estimate uncertainties due to its choice, using Monte Carlo integration [8].

As a typical example, Fig. 3 shows the photon rate for isotropic momentum distribution with $\xi = 0$ as a function of the angle ϕ in the x - y -plane (see Fig. 1), separated in the two circularly polarized states and their unpolarized sum. Results for other ξ values look qualitatively similar. For the QGP where the hard energy scale exceeds $T \gtrsim 200$ MeV, the ratio $E/T = 5$ corresponds to photon energies of 1 GeV or higher. Photons with such high energy are emitted during the initial stage and are not affected by the subsequent expansion. Background effects typically set in at lower energies. This is complementary to the studies in [9] where much lower photon energies were considered. One observes a maximum for left (right) circular polarized photons along (opposite to) the direction of the global spin. This is in accordance with

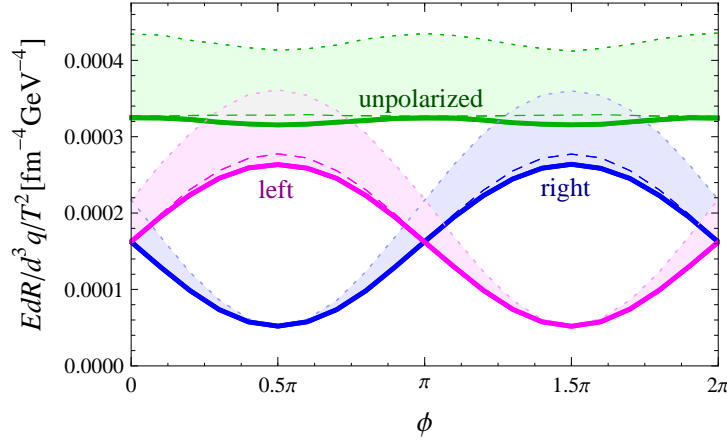


Figure 3. (Color online). Photon rate for $E/T = 5$ and $\xi = 0$ for full quark polarization $\bar{p}_{\text{rest}} = 1$, separated in left and right circularly polarized and unpolarized photons, as a function of the azimuthal angle ϕ in the x - y -plane. The hard-soft separation scale is varied by a factor 2 up (dashed line) or down (dotted line) around its central value (thick line).

high-energy Compton scattering, where helicity states couple to circularly polarized photons. There appears to be some angular dependence also for the unpolarized photon rate [25,26,27]. We verified that this effect is physical for prolate momentum anisotropies with $\xi \approx -1$ (corresponding to early stages of the collision) where it arises from the angular dependence of unpolarized photon emission from quark-antiquark annihilation with transverse spin. For isotropic and oblate systems in momentum space with $\xi \gtrsim 0$, the unpolarized photon rate only shows small ϕ -dependence such that the non-negligible dependence of the unpolarized effect on the choice of the intermediate cutoff scale k^* prevents definitive conclusions at this point. Apart from circular polarization, the remaining Stokes parameters that describe linear polarization all vanish.

From the maximum (max) and minimum (min) photon rate as a function of ϕ we obtain the visibility $V = (\text{max} - \text{min}) / (\text{max} + \text{min})$. Figure 4 depicts V for polarized photons as a function of the photon energy. In this ratio, systematic uncertainties are largely cancelled. We find that the visibility increases with growing oblate anisotropy $\xi \gtrsim 0$ (corresponding to later stages of the collision) and photon energy, but decreases for prolate distributions $\xi < 0$. For a different choice of the prefactor $N(\xi) = \sqrt{1 + \xi}$ for the anisotropy distribution functions used in [24] one finds a similar dependence on ξ , but less pronounced. For energies $E/T \lesssim 4$, the influence of our assumption (1) cannot be neglected, but for larger energies, the result appears to be rather insensitive to the choice of the intermediate cutoff scale k^* . The visibility is proportional to global polarization $\propto \bar{p}_{\text{rest}}$ (e.g. for $\xi = 10$ and global spin polarization $\bar{p}_{\text{rest}} \approx 0.05$, the visibility would be $V \approx 0.04$), and is largest in the x - y -plane. Thus, the quark polarization is imprinted on emitted photons. Therefore pho-

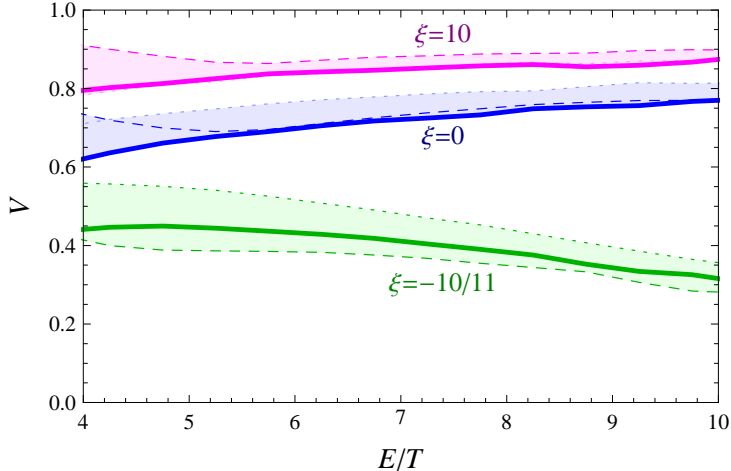


Figure 4. (Color online). Visibility V of photons with circular polarization as a function of photon energy for various anisotropy parameters ξ . Bands mean the same as in Fig. 3.

ton polarization can be used to measure spin polarization, provided one can detect the polarization state of high-energy photons.

Due to the involved dynamics in heavy-ion collisions, there is a rich variety of possible photons sources. One usually separates the contributions into direct photons, fragmentation photons, and background photons [28,29]. Direct photons are those produced through parton collisions, like Compton scattering and quark - antiquark annihilation. Apart from thermal photons from the quark-gluon plasma, direct photons also include prompt photons from initial hard parton scatterings, as well as thermal photons from a later hadron gas phase. Fragmentation photons are produced by bremsstrahlung emitted from final state partons. Background photons are produced by hadron decay in later stages of the collision such as $\pi^0 \rightarrow \gamma\gamma$ or $\eta \rightarrow \gamma\gamma$. Additionally, high-energy partons in the form of jets can produce photons through direct interaction with the medium and through in-medium bremsstrahlung. In the following, we will analyze to what extent these processes could lead to the production of polarized photons.

Prompt photons dominate the photon spectrum at the highest energies. If one considers the leading-order contributions from Compton scattering and annihilation, one sees from Eq. (6) that the photons produced can only be polarized if the incoming quarks are polarized. This could occur through spin polarization of the ions in the storage ring. The polarization of incoming quarks could also be triggered through a strong magnetic field that is created in non-central collisions as explained below. However, this energy range is not of interest for thermal photon production, because at this energy range thermal photons are highly suppressed.

For energies below a few GeV, background photons become increasingly dom-

inating and difficult to subtract [28]. The main contribution of background photons originates from $\pi^0 \rightarrow 2\gamma$ decay. But since the pion has spin zero [30,31], it can not carry any spin information out of the plasma, and thus no net polarization can be transferred to the two photons produced. Thus, background photons from π^0 do not introduce additional photon polarization.

Experimentally, the energy region of a few GeV is the most promising to look for thermal photons. Here photons produced through jets compete with or even exceed the prompt photon yield [29]. Most important are photons that are produced through conversion of a jet in the plasma through Compton scattering or annihilation. Polarized photons could be produced either if the thermal quark from the QGP is polarized or if a quark in the jet is polarized. In the former case a global quark polarization could be reflected in the photons emitted. In the latter case, strong magnetic fields could lead to a polarization of quarks and thus to the emission of polarized photons. It has been pointed out recently that in non-central heavy-ion collisions indeed the charged remnants of the incoming nuclei that do not participate in the collision induce a very strong magnetic field in the collision region that can exceed the critical magnetic field of electrons for a short time [32]. This field would lead to a spin orientation of produced quarks in jets which in turn could be converted into polarized photons through interaction with the hot medium. A strong magnetic field would further allow for magnetic bremsstrahlung [33], that is the emission of a photon from a quark or antiquark in the presence of a background magnetic field - a process that is kinematically forbidden for a vanishing magnetic field. A careful quantitative study of all these different possible contributions is necessary to estimate the experimental detectability of the polarization produced in the quark-gluon plasma.

As unpolarized photons are routinely observed in high-energy collisions [7,25], a main experimental challenge is to detect circularly polarized photons in the GeV energy range [34,35]. We propose to use a single aligned crystal acting as a quarter wave plate for high energy photons [36,37] to convert their circular polarization into linear polarization that can then be analyzed by means of another aligned crystal through coherent electron-positron pair production [36,35,38]. The low photon fluxes of the order of a photon per collision that are expected from the QGP in the GeV energy range would pose an additional technical difficulty for the experimental detectability. On the other hand, even if the polarization detector covers only a fraction of the 4π coverage at some fixed position in the laboratory, angular dependence studies can be performed as the reconstructible reaction plane and thus the global polarization changes its orientation from event to event.

Concluding, we have shown that spin polarization in the QGP would lead to the emission of circularly polarized photons. Good photon polarization visibility is found in particular for higher photon energies and oblate momentum

distributions. We further showed that for the leading order processes considered, momentum anisotropy alone cannot give rise to polarization of photons. With the ongoing experimental progress in the detection of polarized photons in the GeV range, further exploration of other sources of polarized photons in heavy-ion collisions would certainly be desirable.

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